

CONSTITUENT STRING MODEL FOR HYBRID MESONIC EXCITATIONS

Yu.S.Kalashnikova¹, Yu.B.Yufryakov²

Institute of Theoretical and Experimental Physics
117259, Russia, Moscow, B.Chermushkinskaya 25.

Abstract

The model for hybrid excitations of the QCD string with quarks is presented starting from the perturbation theory in the nonperturbative background. The propagation of a system containing $q\bar{q}$ -pair and gluon is considered. The simplified version of the Hamiltonian, including both long-range nonperturbative interaction and Coulomb force, is derived. The masses of the lowest $q\bar{q}g$ hybrids are evaluated, and numerical results for the spectra are listed.

1 Introduction

One of the most important features of the QCD Lagrangian is the presence of gluonic degrees of freedom which should exhibit themselves at the constituent

¹E-mail: yulia@vxitep.itep.ru

²E-mail: yufryakov@vxitep.itep.ru

level, namely in the form of glueballs and hybrids, i.e. non- $q\bar{q}$ exotics. Unfortunately, the current experimental situation is too complicated to give us unambiguous proof that such states exist; some candidates for non- $q\bar{q}$ exotics appear and disappear from time to time, changing their masses and quantum numbers (for the up-to-date review see [1]). Nevertheless, there is no doubts that standard $q\bar{q}$ nonets are overpopulated, but the question about the nature of the "extra" states is still open. On the other hand, there is no substantial progress in the description of the strong coupling nonperturbative regime in QCD. There are some QCD – inspired theoretical approaches, but none of them are able to provide reliable enough predictions for masses and decay rates of hybrid mesons.

The QCD sum rules estimations for exotic hybrids are rather unstable: first results [2] predicted for the 1^{-+} light hybrid the mass 1.2-1.7 GeV, while more recent calculations give 2.1 GeV [3] and 2.5 GeV [4].

In the bag model the gluons are automatically transverse, and the lowest electric gluon (with $J^P = 1^+$) in the spherical cavity is much lighter than the lowest magnetic one (with $J^P = 1^-$). The lowest hybrids with light quarks have the mass about 1.5 GeV [5,6]. The masses of hybrids with heavy quarks were estimated in [7] taking into account the bag deformation, with the results 3.9 GeV for $c\bar{c}$ hybrid and 10.5 GeV for $b\bar{b}$ one.

Constituent gluon model was introduced in [8,9]. In this model the linear potential is introduced *ad hoc*, in analogy with the charmonium system. It was proposed also that the gluon orbital momentum is diagonal. As a consequence, in such model the lowest states with the S -wave gluon have non-exotic quantum numbers. The mass predictions of constituent model are 1.3-1.8 GeV for the lowest non-exotic hybrids.

Flux-tube model [10] predicts degenerate (up to spin corrections) lightest hybrid states at 1.8-1.9 GeV. In this model phonon-type excitations of the string connecting quark and antiquark are interpreted as hybrids. First calculations [10] assumed small oscillation approximation, and recent results [11] demonstrate that this approximation might be inadequate. Improved version [11] is given in the framework of "one-bead" flux-tube model.

Here we present the studies of the $q\bar{q}g$ system in the framework of Vacuum Background Correlator method [12]. The main assumption is that nonperturbative background fields $\{B_\mu\}$ exist in the QCD vacuum, which ensure

the area law asymptotic for the Wilson loop along the closed contour C ,

$$\langle W(C) \rangle_B \rightarrow N_c \exp(-\sigma S),$$

with $S = S_{min}$ being the minimal surface inside the contour C . It was confirmed by cluster expansion method [12] that if one assumes the existence of finite correlation length for the background, the asymptotical behaviour of the Wilson loop average is compatible with the area law. The deviation from the area law at large distances with $S = S_{min}$ are caused by the perturbations over the background. The "minimal" area law gives the string-type interaction in the $q\bar{q}$ system, while the perturbative fields are responsible for the string vibrations [13].

2 Green function for the $q\bar{q}g$ system in the Vacuum Background Correlator method

The constituent gluon in the Vacuum Background Correlator method is introduced as a gluon propagating in the nonperturbative background field [14,15]. Following [15], we split the gluonic field A_μ into the background field B_μ and the perturbation a_μ over the background.

$$A_\mu = B_\mu + a_\mu \quad (1)$$

We ascribe the inhomogeneous part of gauge transformation to the field B_μ ,

$$B_\mu \rightarrow U^+(B_\mu + \frac{i}{g}\partial_\mu)U, \quad a_\mu \rightarrow U^+a_\mu U, \quad (2)$$

so that the states involving the field a_μ may be formed in the gauge-invariant manner. One-gluon hybrid is represented as

$$\Psi(x_q, x_{\bar{q}}, x_g) = \bar{\psi}_\alpha(x_{\bar{q}})\Phi_\beta^\alpha(x_{\bar{q}}, x_g)a_\gamma^\beta(x_g)\Phi_\delta^\gamma(x_g, x_q)\psi^\delta(x_q), \quad (3)$$

where $\alpha\dots\delta$ are the colour indices in the fundamental representation, $a_\gamma^\beta = a_a(\lambda_a)_\gamma^\beta$, and parallel transporters Φ contain only background field:

$$\Phi_\beta^\alpha(x, y) = (P \exp \int_y^x B_\mu dz_\mu)_\beta^\alpha \quad (4)$$

The Green function for the $q\bar{q}g$ system is obtained by averaging the product $\Psi_{in}\Psi_{out}^+$ over the background field configurations:

$$G(x_q x_{\bar{q}} x_g, y_q y_{\bar{q}} y_g) = \langle \Psi_{in}(y_q, y_{\bar{q}}, y_g) \Psi_{out}^+(x_q x_{\bar{q}} x_g) \rangle_B. \quad (5)$$

The dynamics of the field a_μ is defined, in accordance with the decomposition (1), by expanding the QCD Lagrangian up to the second order in the fields a_μ (in the Euclidean space-time)

$$\begin{aligned} L(a) = & -\frac{1}{4}(F_{\mu\nu}(B))^2 + a_\nu D_\mu(B) F_{\mu\nu}(B) + \\ & + \frac{1}{2} a_\nu (D_\lambda D_\lambda \delta_{\mu\nu} - D_\mu D_\nu - g F_{\mu\nu}(B)) a_\mu, \\ & D_\lambda^{ca} = \partial_\lambda \delta^{ca} + g f^{cba} B_\lambda^b, \end{aligned} \quad (6)$$

with the background gauge fixing term

$$G^a = \partial_\mu a_\mu^a + g f^{abc} B_\mu^b a_\mu^c \quad (7)$$

In what follows we skip the issue of ghosts. The linear in the fields a_μ part of the Lagrangian (6) disappears if the field B_μ satisfies the classical equation of motion $D_\mu F_{\mu\nu} = 0$. To be on the safe side one is to assume that the background is the classical one, or at least, that the transition vertex generated by this term is small.

The Green function for the field a_μ propagating in the given background B_μ may be identified in the background gauge as

$$G_{\mu\nu}^{-1} = D^2 \delta_{\mu\nu} - D_\nu D_\mu - g F_{\mu\nu} + \frac{1}{\xi} D_\mu D_\nu = M_{\mu\nu} + \frac{1}{\xi} D_\mu D_\nu. \quad (8)$$

If the classical equations of motion are satisfied, then one has $M_{\mu\nu} D_\nu = 0$, and the Green function (8) may be rewritten as

$$G_{\mu\nu} = (\delta_{\mu\lambda} + (\xi - 1) D_\mu \frac{1}{D^2} D_\lambda) (D^2 - 2gF)_{\lambda\nu}^{-1}. \quad (9)$$

The choice $\xi = 0$ corresponds to the Landau gauge, in which the Green function (9) contains explicitly the projector $P_{\mu\lambda}$ onto transverse states:

$$P_{\mu\lambda} = \delta_{\mu\lambda} - D_\mu \left(\frac{1}{D^2} \right) D_\lambda. \quad (10)$$

To define the effective action for the $q\bar{q}g$ system we use the Feynman–Schwinger representation [16]. To do it in proper way one should take into account spin degrees of freedom of quarks and gluon. Here we omit the spin dependence, reducing the problem to the scalar one. This simplified version of the model corresponds to the neglecting of colour magnetic interaction (the term proportional to the gF in (9)), and omitting the projector. Similarly, spin dependence in the quark Green function is also omitted, and we assume

$$G_q = (D^2 - m_q^2)^{-1}. \quad (11)$$

As the result, the Feynman–Schwinger representation for the hybrid Green function takes the form

$$G(x_q x_{\bar{q}} x_g, y_q y_{\bar{q}} y_g) = \int_0^\infty ds \int_0^\infty d\bar{s} \int_0^\infty d\bar{S} \int Dz D\bar{z} DZ \exp(-\mathcal{K}_q - \mathcal{K}_{\bar{q}} - \mathcal{K}_g) < \mathcal{W} >_B, \quad (12)$$

where

$$\mathcal{K}_q = m_q^2 s + \frac{1}{4} \int_0^s \dot{z}^2(\tau) d\tau, \quad \mathcal{K}_{\bar{q}} = m_{\bar{q}}^2 \bar{s} + \frac{1}{4} \int_0^{\bar{s}} \dot{\bar{z}}^2(\tau) d\tau, \\ \mathcal{K}_g = \frac{1}{4} \int_0^S \dot{Z}^2(\tau) d\tau,$$

with boundary conditions

$$z(0) = y_q, \quad \bar{z}(0) = y_{\bar{q}}, \quad Z(0) = y_g, \\ z(s) = x_q, \quad \bar{z}(\bar{s}) = x_{\bar{q}}, \quad Z(S) = x_g,$$

and \mathcal{W} is the Wilson loop operator

$$\mathcal{W} = (\lambda_a)_\beta^\alpha (\Phi_{\Gamma_q}(y_q, x_q))_\gamma^\beta (\lambda_b)_\delta^\gamma (\Phi_{\Gamma_{\bar{q}}}(x_{\bar{q}}, y_{\bar{q}}))_\alpha^\delta (\Phi_{\Gamma_g}(y_g, x_g))_{ab}, \quad (13)$$

which corresponds to the propagation of quark along the path Γ_q , of antiquark along the path $\Gamma_{\bar{q}}$ and of gluon along the path Γ_g (see Fig.); here a, b are the colour indices in the adjoint representation. As in the case of $q\bar{q}$ system [17], all the dependence on the background field is contained in the Wilson loop operator (13).

The Wilson loop configuration (13) may be rewritten using the relation between ordered exponents along the gluon path Γ_g in the adjoint and fundamental representations:

$$\frac{1}{2}(\Phi_{\Gamma_g}(x, y))_{ab} = (\lambda_a)_\beta^\alpha (\Phi_{\Gamma_g}(x, y))_\gamma^\beta (\lambda_b)_\delta^\gamma (\Phi_{\bar{\Gamma}_g}(y, x))_\alpha^\delta, \quad (14)$$

where the path $\bar{\Gamma}_g$ coincides with the path Γ_g and is directed oppositely. The result reads

$$\mathcal{W} = \frac{1}{2} Sp W_1 Sp W_2 - \frac{1}{2N_c} Sp W, \quad (15)$$

where W_1, W_2 and W are the Wilson loops in the fundamental representation along the closed contours $C_1 = \Gamma_q \bar{\Gamma}_g$, $C_2 = \Gamma_{\bar{q}} \Gamma_g$ and $C = \Gamma_q \bar{\Gamma}_{\bar{q}}$ shown at the Figure.

3 Generalized area law and effective Hamiltonian

To average the Wilson loop configuration (15) over the background we use the cluster expansion method generalized in [18] to consider the average of more than one Wilson loop. For the contours C_1 and C_2 with the average size much larger than the gluonic correlation length T_g we arrive to the generalized area law

$$\langle \mathcal{W} \rangle = \frac{N_c^2 - 1}{2} \exp(-\sigma(S_1 + S_2)), \quad (16)$$

where σ is the string tension in the fundamental representation, and S_1 and S_2 are the minimal surfaces inside the contours C_1 and C_2 . The area law (16) holds for all the configurations in the $q\bar{q}g$ system, apart from the special case of the contours C_1 and C_2 embedded into the same plane, where, instead of (16), one has

$$\langle \mathcal{W} \rangle = \frac{N_c^2 - 1}{2} \exp(-\sigma(S_1 - S_2)) - \sigma^{adj} S_2, \quad S_1 > S_2, \quad (17)$$

σ^{adj} is the string tension in the adjoint representation. The regimes (16) and (17) match smoothly each other at the distances between the contours C_1 and C_2 of order of correlation length T_g .

If the string tension is defined mainly by the contribution of second order correlators, then $\sigma^{adj}/\sigma = 9/4$ for the $SU(3)$ colour group. On the other hand, the area law for the Wilson loop in the adjoint representation was observed on the lattice [19] with $\sigma^{adj}/\sigma \approx 2$, and the same result holds true in the limit $N_c \rightarrow \infty$. Having all this in mind, we assume the regime (16) to be valid everywhere in the $q\bar{q}g$ configuration space.

The four-dimensional dynamics in (12) can be reduced to the three-dimensional one following the procedure used in [17]. Namely, choosing the physical time parametrization

$$z_\mu = (\tau, \vec{r}_q), \quad \bar{z}_\mu = (\tau, \vec{r}_{\bar{q}}), \quad Z_\mu = (\tau, \vec{r}_g)$$

and introducing new dynamical variables

$$\mu_1(\tau) = \frac{T}{2s} \dot{z}_0(\tau), \quad \mu_2(\tau) = \frac{T}{2\bar{s}} \dot{\bar{z}}_0(\tau), \quad \mu_3(\tau) = \frac{T}{2S} \dot{Z}_0(\tau),$$

with $0 \leq \tau \leq T$, we arrive to the three-dimensional representation for the Green function:

$$G = \int D\vec{r}_q D\vec{r}_{\bar{q}} D\vec{r}_g D\mu_1 D\mu_2 D\mu_3 \exp(-A), \quad (18)$$

with the effective action

$$A = \int_0^T d\tau \left\{ \frac{m_q^2}{2\mu_1} + \frac{m_{\bar{q}}^2}{2\mu_2} + \frac{\mu_1 + \mu_2 + \mu_3}{2} + \frac{\mu_1 \dot{r}_q^2}{2} + \frac{\mu_2 \dot{r}_{\bar{q}}^2}{2} + \frac{\mu_3 \dot{r}_g^2}{2} + \right. \quad (19)$$

$$\left. + \sigma \int_0^1 d\beta_1 \sqrt{\dot{w}_1 w_1'^2 - (\dot{w}_1 w_1')^2} + \sigma \int_0^1 d\beta_2 \sqrt{\dot{w}_2 w_2'^2 - (\dot{w}_2 w_2')^2} \right\},$$

where the surfaces S_1 and S_2 are parametrized by the coordinates $w_{i\mu}$, $\dot{w}_{i\mu} = \frac{\partial w_{i\mu}}{\partial \tau}$, $w'_{i\mu} = \frac{\partial w_{i\mu}}{\partial \beta_i}$, $i = 1, 2$. Assuming the straight-line ansatz for the minimal surfaces,

$$w_{1\mu} = \beta_1 z_\mu + (1 - \beta_1) Z_\mu, \quad w_{2\mu} = \beta_2 \bar{z}_\mu + (1 - \beta_2) Z_\mu,$$

we write out the effective Lagrangian for the $q\bar{q}g$ system as

$$L = \frac{m_q^2}{2\mu_1} + \frac{m_{\bar{q}}^2}{2\mu_2} + \frac{\mu_1 + \mu_2 + \mu_3}{2} + \frac{\mu_1 \dot{r}_q^2}{2} + \frac{\mu_2 \dot{r}_{\bar{q}}^2}{2} + \frac{\mu_3 \dot{r}_g^2}{2} + \quad (20)$$

$$\begin{aligned}
& +\sigma\rho_1\int_0^1 d\beta_1\sqrt{1+l_1^2}+\sigma\rho_2\int_0^1 d\beta_2\sqrt{1+l_2^2}, \\
& \vec{l}_1=\frac{1}{\rho_1}\vec{\rho}_1\times(\beta_1\dot{\vec{r}}_q+(1-\beta_1)\dot{\vec{r}}_g), \\
& \vec{l}_2=\frac{1}{\rho_2}\vec{\rho}_2\times(\beta_2\dot{\vec{r}}_{\bar{q}}+(1-\beta_2)\dot{\vec{r}}_g), \\
& \vec{\rho}_1=\vec{r}_q-\vec{r}_g, \quad \vec{\rho}_2=\vec{r}_{\bar{q}}-\vec{r}_g.
\end{aligned}$$

To obtain the effective Hamiltonian one should define the momenta and express the velocities in terms of momenta. It cannot be done explicitly because of presence of square roots in (20), and to deal with this problem the auxiliary field approach was suggested in [17]. However, it was shown in [17] that for the low values of relative orbital momenta the square roots in (20) can be expanded up to the second order in the angular velocities \vec{l}_i , the approximation proved to be accurate enough even for the massless constituents. Within this approximation the problem is reduced to the potential – like one, while the terms $\sim l_i^2$ can be taken into account perturbatively. The corresponding Hamiltonian in the Minkowsky space–time in the centre–of–mass frame is easily obtained from (20):

$$H_0 = \frac{m_q^2}{2\mu_1} + \frac{m_{\bar{q}}^2}{2\mu_2} + \frac{\mu_1 + \mu_2 + \mu_3}{2} + \frac{p^2}{2\mu_p} + \frac{Q^2}{2\mu_Q} + \sigma\rho_1 + \sigma\rho_2, \quad (21)$$

$$\vec{\rho}_1 = \vec{\rho} - \frac{\mu_2}{\mu_1 + \mu_2}\vec{r}, \quad \vec{\rho}_2 = -\vec{\rho} - \frac{\mu_1}{\mu_1 + \mu_2}\vec{r},$$

where the Jacobi coordinates

$$\vec{r} = \vec{r}_1 - \vec{r}_2, \quad \vec{\rho}_2 = \vec{r}_3 - \frac{\mu_1\vec{r}_1 + \mu_2\vec{r}_2}{\mu_1 + \mu_2},$$

and conjugated momenta \vec{p} and \vec{Q} are introduced, and μ_p and μ_Q are the reduced masses

$$\mu_p = \frac{\mu_1\mu_2}{\mu_1 + \mu_2}, \quad \mu_Q = \frac{\mu_2(\mu_1 + \mu_2)}{\mu_1 + \mu_2 + \mu_3}.$$

The Hamiltonian still contains the fields $\mu_i(\tau)$, and the integration over $\{\mu_i\}$ is to be performed in the path integral representation (18) (or, equivalently, taking the extremal values of μ_i in the Hamiltonian). Only after that the quantization should be carried out.

Technically, it is more convenient to proceed in a way suggested in [20]: first find the eigenvalues of the Hamiltonian (21) assuming μ_i to be c -numbers, and after that minimize the eigenenergies in μ_i . This procedure works with rather good accuracy for the lowest states, and reduces the problem to the nonrelativistic three-body one, with μ_i playing the role of constituent masses. We note that although the Hamiltonian (21) looks like the Hamiltonian of the nonrelativistic potential model, it is essentially relativistic, and the masses μ_i are not introduced by hand, but are calculated and expressed in terms of string tension σ and quark masses.

Another advantage of the above-described method is that it allows for the approximate solution to the problem of separating out the physical transverse states. Indeed, let us impose the constraint

$$\mu_3 \Psi_0 - \mu_3 (\dot{\vec{r}}_g \vec{\Psi}) = 0 \quad (22)$$

to project out the physical hybrid state $\Psi_\lambda = (\Psi_0, \vec{\Psi})$, where λ is the gluon spin index. The constraint (22) is compatible with the projector $P_{\mu\lambda}$ (10) after averaging over background and introducing the variables μ_i . In the potential-like regime one has $\vec{p}_g = \mu_3 \dot{\vec{r}}_g$, and we choose the physical states to be transverse with respect to the three-dimensional gluon momentum:

$$\vec{p}_3 \vec{\Psi} = 0, \quad \Psi_0 = 0. \quad (23)$$

We are forced to impose the condition (23), because we have neglected the spin dependence in the gluon Green function (9), so the condition (23) should be treated as variational ansatz. The rigorous analysis of the transverse and longitudinal gluonic degrees of freedom should be done with the inclusion of spin into the path integral representation for the gluon Green function.

4 Numerical results and discussion

In the actual calculations the Hamiltonian was supplied with the short range Coulomb interaction

$$V_c = \frac{\alpha_s}{6r} - \frac{3\alpha_s}{2\rho_1} - \frac{3\alpha_s}{2\rho_2}. \quad (24)$$

As a variational ansatz the Gaussian type radial wave functions were chosen. The constraint (23) was satisfied by taking the orbital wave functions

diagonal in the total angular momentum j in the gluonic subsystem, so that the states contain electric or magnetic gluon:

$$\vec{\Psi}_j^e \sim \vec{Y}_{jjm}(\hat{Q}), \quad \vec{\Psi}_j^m \sim \sqrt{\frac{j+1}{2j+1}} \vec{Y}_{jj-1m}(\hat{Q}) + \sqrt{\frac{j}{2j+1}} \vec{Y}_{jj+1m}(\hat{Q}). \quad (25)$$

With this choice the electric and magnetic hybrids are degenerate, and this degeneracy will be removed by string corrections and by spin-dependent force.

The quantum numbers of a one-gluon hybrid are given by

$$P = (-1)^{l+j}, \quad C = (-1)^{l+s+1} \quad (26)$$

for the states with electric gluon, and

$$P = (-1)^{l+j+1}, \quad C = (-1)^{l+s+1} \quad (27)$$

for the states with magnetic gluon, where l and s are the angular momentum and total spin in the quark-antiquark subsystem. So the possible quantum numbers for the ground state are

$$J^{PC} = 0^{\mp+}, 1^{\mp+}, 2^{\mp+}, 1^{\mp-}, \quad (28)$$

where the upper/lower sign stands for the state with electric/magnetic gluon (26)/(27).

The most complicated problem in the constituent approaches is not the relative arrangement of ground and excited states, but the absolute scale of masses. In the potential model large negative constant is needed to fit the $q\bar{q}$ spectrum, and this constant is different for the sectors with different flavour content. In the described approach the perimeter terms for the Wilson loop and/or hadronic shifts might be responsible for the constant term. We use the prescription that for the hybrid state with two strings the additive constant is twice as large as for the $q\bar{q}$ meson with only one string.

So, the procedure used is: 1) to define the constant term for the given values of parameters from the fit to the S -, P - and D - wave meson levels, and 2) to calculate the hybrid mass with the constant multiplied by two.

The numerical results for the spectra of hybrids with light quarks are listed in the Table 1 for different values of quark mass, string tension and α_s .

Slightly another procedure was used to calculate the masses of hybrids with heavy quarks: the constant term was taken from the fit to heavy–light ($D-$ and $B-$) meson masses. We think that the ”heavy–light” constants are more consistent phenomenologically for the hybrid with two ”heavy–light” strings. The results for the ground states are given in Table 2.

The spectra we have obtained are rather similar to the ones of the flux–tube model [10,11]. There is a lot of common in these two approaches, because both models try to account for string vibrations. Moreover, the Hamiltonian (21) looks quite similar to the Hamiltonian of the ”one–bead” flux–tube model [11]. As it was already mentioned, the numerical calculations [11] do not support the small oscillation approximation proposed in the original version [10] of the flux–tube model. Hence, the constraint imposed that the ”beads” can oscillate only in the transverse (with respect to quark–antiquark) direction seems to be a little suspicious. However, the heavy quark hybrid system was analysed [21] in the string–type regime of the Lagrangian (20), which matches smoothly at low j the potential–type regime described here, and it was shown that the effective values of $\sigma\rho_1$ and $\sigma\rho_2$ are equal, and it is just the case of the ”one–bead” flux–tube model. The difference comes from the fact that in the flux–tube the masses of constituents (including the bead) are fixed, while in our approach the effective masses are the variables. Another numerical discrepancy is due to the constant term: in the flux–tube there is one string, so it is reasonable to use the constant fitted by the $q\bar{q}$ spectrum, while we have two distinguishable strings, and the constant defined from the $q\bar{q}$ spectrum should be multiplied by the factor of two. These discrepancies compensate each other, so the almost exact coincidence of the results seems to be to some extent accidental.

The most important difference between the models is in quantum numbers. In the presented picture the confinement is of the stochastic nature and is ensured by the background fields, with well–defined perturbation theory in the background. As a result, the constituent gluon carries quantum numbers of its own. On the contrary, the flux–tube, being motivated by the strong coupling expansion, knows nothing about the gluons which populate the string. The excitations of the string are described by the collective phonon–type modes, so instead of (28) one has for the ground state

$$J^{PC} = 0^{\mp\pm}, 1^{\mp\pm}, 2^{\mp\pm}, 1^{\mp\mp}. \quad (29)$$

The most clear decay signature of hybrids is the suppression of a hybrid

decay into two S -wave ground state mesons. This signature takes place for the flux-tube hybrid [22] as well as for the electric constituent hybrid [23], and follows from the symmetry of the wave functions involved. It means that P -odd hybrids (28) and (29) have the same decay properties, and the discrepancy between the models should reveal itself in the P -even hybrid sector. However, all currently discussed hybrid candidates in the light quark sector are P -odd!!

There is growing evidence that hybrids are found at last, the belief based mainly on the above-mentioned signature. Indeed, the $0^{-+}\pi(1800)$ state seen by VES [24] decays mainly into πf_0 with $\pi\rho$ mode suppressed; the exotic 1^{-+} signal is seen in BNL [25] in the πf_1 final state; the $\pi(1775)$ is seen in charge exchange photoproduction [26] decaying into πf_2 , that might be 2^{-+} ; the $\rho'(1460)$ decays mainly into πa_1 in contrast to $2\ ^3S_1\ q\bar{q}$ assignment [27]; the rather promising isoscalar $2^{-+}\eta_2(1870)$ [28] is observed with πa_2 and $\eta f_0(980)$ decay models. Unfortunately, up to now no hybrid-like activity is observed in the P -even sector.

5 Concluding remarks

We have demonstrated that the perturbation theory in the nonperturbative confining background supports the existence of $q\bar{q}g$ bound states. The hybrid mesonic excitation looks like a system of a gluon with two straight-line strings with quarks at the ends. For low values of gluon orbital momentum the problem is reduced to the potential-like one, and the resulting hybrid spectra are compatible with the data on light quark meson spectroscopy.

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Table 1. Predicted masses of hybrids with light quarks

m_q, GeV	α_s	σ, GeV^2	$M_{q\bar{q}g}, GeV$		
			$j = 1, l = 0$	$j = 1, l = 1$	$j = 2, l = 0$
0	0.3	0.18	1.73	2.03	2.2
		0.2	1.68	2.00	2.18
0.1	0.3	0.18	1.71	2.02	2.19
		0.2	1.68	2.00	2.18
0	0.7	0.18	1.60	1.95	2.15
0.1	0.7	0.18	1.58	1.93	2.14

Table 2. Predicted masses of hybrids with heavy quarks;
 $\sigma = 0.18 GeV^2, \alpha_s = 0.3$

m_c, GeV	$M_{c\bar{c}g}, GeV$
1.2	4.12
1.5	4.11
1.7	4.10
m_b, GeV	$M_{b\bar{b}g}, GeV$
4.5	10.64
5.2	10.64

Figure caption

Wilson loop configuration corresponding to the propagation of the hybrid state.